Theoretical Notes
Note 368

JUSTIFICATION AND VERIFICATION OF HIGH-ALTITUDE EMP THEORY

PART I

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SECTION 1
INTRODUCTION

Over the 22 years since the first publication (Ref. 1) of the theory of High-Altitude Electromagnetic Pulse (HEMP), there have been several doubters of the correctness of that theory. On one occasion it was briefly claimed that the HEMP is a much larger pulse than our theory indicates, and is a longitudinal wave rather than transverse. This claim was easily shown to rest on an incorrect application of a standard formula for the fields of a charge moving at relativistic speed. More commonly, it has been claimed that the HEMP is a much smaller pulse than our theory indicates and it has been implied, though not directly stated in writing, that the HEMP has been exaggerated by those who work on it in order to perpetuate their own employment. It could be noted that, in some quarters, the disparagement of HEMP has itself become an occupation. While we have found that no amount of technical reasoning suffices to quiet such criticism, we have learned to live with it, and even to regard it as possibly having some beneficial effects, for example in bringing the question of the HEMP threat to electrical and electronic systems to the attenuation of a wider circle of individuals who have responsibility for those systems.

Thus our principal concern is to convince individuals, with technical backgrounds and open minds, who for various reasons have not been convinced by previously published papers on HEMP theory, that the theory and calculated results are at least approximately correct. One possible difficulty with previous papers (Refs. 1, 2, 3) is that they are based on solving Maxwell's equations. While that is the most legitimate approach for the mathematically inclined reader, many of the individuals we think it important to reach may not feel comfortable with that approach. We admit to being surprised at the number of people who have wanted to understand HEMP in terms of the fields radiated by individual Compton recoil electrons. Apparently our schools do a better job in teaching the applications of Maxwell's equations (in this case, the cyclotron radiation) than they do in imparting a basic understanding of those equations and how they work. However, the generation of the HEMP can be understood quite well on the basis of concepts familiar to all electrical engineers and physicists. A derivation of the outgoing wave equation based on these concepts is presented in Section 2. This equation, though extremely simple, gives the HEMP to an accuracy of a few percent in practical cases. We also show that adding the fields radiated by individual electrons gives exactly the same answer in a simple but relevant example.

The confidence we have in our calculations of the HEMP rests on two circumstances. The first of these is the basic simplicity of the theory. The physical processes involved, e.g., Compton scattering, are quite well known, and the physical parameters needed in the calculations, such as electron mobility, have been measured in relevant laboratory experiments. There is no mathematical difficulty in determining the solution of the outgoing wave equation, or in understanding why it is an accurate approximation. Nevertheless, to acquire a sufficient understanding of HEMP to be able to say, of one's own knowledge, that our answers are right to 10 or 20% (or wrong) is not a trivial exercise. While the concepts we start with are familiar, in applying them we shall soon come into new ground for many readers. One will not find our problem worked out in textbooks. For example, the model of cyclotron radiation from individual Compton recoil electrons is very difficult to apply with accuracy to our problem because of the multitudinous secondary electrons, which absorb the radiation emitted by the Compton electrons. Readers who have the patience to follow the development in Section 2 will see that the outgoing wave theory provides a way to avoid unnecessary complications and reduces the problem to its barest essentials.

The other circumstance is that there is experimental data on the HEMP obtained by the Los Alamos Scientific Laboratory in the nuclear test series carried out in 1962. In a classified companion report (Ref. 4) we present calculations of the HEMP from the Kingfish and Bluegill events and compare them with the experimental data. These calculations were performed some years ago, but have not been widely circulated. In order to make the calculations transparently honest, the gamma-ray output was provided by Los Alamos, the HEMP calculations were performed by MRC and the comparison with the experimental data was made by RDA. The degree of agreement between calculation and experiment gives important verification of the correctness of HEMP theory.

A feature of HEMP theory that has troubled some people is that it determines the radiated fields by integrating along the single ray from nuclear burst to observer. The angular derivatives in Maxwell's equations are dropped, which amounts to neglecting diffraction in the propagation of the HEMP. The first paper on HEMP (Ref. 1) contained a justification of this neglection, based on the large ratio of the transverse dimension of the radiating volume compared with the wavelengths in the HEMP. However, we have always wondered just how large a correction diffraction would make for Kingfish, a fairly severe geometry. In Section 3, a method for calculating the effect of diffraction is developed. The method is applied in Reference 4 to that event and the effect is found to be less than 1%. This calculation starts from Maxwell's equations, but we hope that many readers will follow the not-very-complicated mathematics.

SECTION 2
THEORY OF THE GENERATION OF HEMP

2.1 GAMMA RAYS AND COMPTON SCATTERING

Nuclear bombs emit a small fraction, of the order of 0.003, of their energy in gamma rays. Thus a 1-megaton bomb, which produces total energy of about 4.2×10^{15} J (Joules) may emit 3 kilotons or about 1.2×10^{13} J in gamma rays. Gamma rays are electromagnetic waves, like radio waves or visible light, but of much higher frequency than either of these--in the range of 10^{20} to 10^{21} Hz. They travel at the speed of light ($c = 3\times10^{8}$ m/sec), and so have wavelengths of the order of 10^{-10} cm. This is smaller than the diameter of atoms (~10^{-8} cm), and in interacting with atoms, gamma rays act more like particles than waves. As a particle, or quantum of electromagnetic radiation, a gamma ray has an energy of the order

$$E_{\gamma} \approx 2 \text{ MeV} \approx 3.2\times10^{-13} \text{ J} \quad . \tag{1}$$

(The MeV unit of energy is a convenient one in nuclear physics; it is equal to the energy gained by an electron in falling through a potential drop of 10^{6} V.) Thus the total number of gammas emitted by a 1 megaton bomb may be of the order

$$N_{\gamma} \approx 4\times10^{25} \text{ gammas} \quad . \tag{2}$$

The principal interaction of gamma rays with air atoms, or other matter, is Compton scattering. In this process, the gamma collides with an electron in the air atom and knocks it out of the atom. In so doing, the gamma transfers part of its energy (on the average about half) to the electron, and is scattered into a new direction. The Compton recoil electron goes generally near the forward direction of the original gamma, never in backward directions. Thus a directed flux of gammas produces a directed electric current of Compton recoil electrons. This current produces the EMP.

The mean free path for Compton scattering in sea-level air is of the order

$$\lambda_\gamma \approx 180 \text{ m} \quad \text{(at sea level)} \quad . \tag{3}$$

At higher altitudes, where the air density is smaller, λ_γ is correspondingly longer. Since the total mass of air in the atmosphere is about 1000 gm/cm^2, and since the mean free path can also be expressed as

$$\lambda_\gamma = 22 \text{ gm/cm}^2 \quad \text{(at any altitude)} \quad , \tag{4}$$

the total number of mean free paths in the atmosphere for gammas coming vertically downwards is

$$N_\lambda = 1000/22 \approx 45 \quad . \tag{5}$$

Only a very small fraction, $\exp(-N_\lambda) \approx 10^{-20}$, of the gammas reach the ground without being scattered. Most of the gammas will suffer their first scattering near the altitude where they have traversed one mean free path. This is the altitude above which there are 22 gm/cm^2 of air and is about

$$z_s = 30 \text{ km} \quad . \tag{6}$$

Because the air density ρ_a falls approximately exponentially with altitude z,

$$\rho_a \approx \rho_0 \exp(-z/h) \quad , \tag{7}$$

where the scale height h is

$$h \approx 6.7 \text{ km} \quad \text{at 30 km altitude} \quad , \tag{8}$$

the local mean free path at this altitude is approximately equal to h. Thus most of the gammas make their first scattering in a layer of air of thickness h centered about z_s. This is the source region of the HEMP. Scattered gammas make little further contribution to the HEMP, for reasons to be explained below.

We can now calculate the density of Compton electrons in the source region. Let us put our nominal 1-megaton bomb at an altitude of 100 km, 70 km above the source region. Then the number of gammas incident on the source region per unit area, assuming they are emitted isotropically by the bomb, is

$$N_A = N_\gamma/4\pi r^2 = 6.5\times10^{14} \text{ gammas/m}^2 \quad . \tag{9}$$

Since every gamma will make a Compton electron, the number of first-scatter Compton electrons per unit area is also N_A. On the assumption that the gammas are scattered in a height h, the number of birth places of first-scatter Compton electrons per unit volume is

$$N_{bp} = N_A/h = 1.0\times10^{11} \text{ b.p.'s/m}^3 \quad . \tag{10}$$

The density of Compton electrons would be the same as this number if they did not move. Because they move in the same direction as the gammas with an average speed of about 0.9 c, their actual density is about 10 times this number shortly after birth, or about $10^{12}/m^3$. (This point is explained in Section 2.2 below.)

An important question is whether the density of Compton electrons is large enough to justify regarding the current they make as a continuous function. The answer to this question depends on the wavelength of electromagnetic fields that they produce and that are of interest. The duration of the HEMP (or that part of it that is of interest to us) is at least 10^{-8} second, so that wavelengths are a few meters. If we wish to resolve the Compton current on the scale of one quarter of a wavelength, or about 1 meter, we need to consider the number of Compton electrons in a volume of 1 m^3. This is about 10^{12} electrons. The average fluctuation in this number is about $\sqrt{10^{12}} = 10^6$, or 1 part in 10^6. Thus it seems clear that, for the wavelengths of interest, the Compton current density can be quite accurately treated as a continuous function. The coherence of the radiated fields is discussed further in Section 2.3 below.

Note that if the Compton electrons all moved together at nearly the speed of light, they would make a current density

$$J_c \approx 48 \text{ A/m}^2 \quad , \qquad (11)$$

a substantial current density.

2.2 THE MOTION OF COMPTON RECOIL ELECTRONS

For 2-MeV gammas, the Compton recoil electrons have an average kinetic energy of about 1 MeV. The most energetic electrons move in the forward direction of the original gamma with kinetic energy 1.78 MeV. The angular distribution of the electrons (per unit solid angle) peaks in the forward direction, and has half of its peak value at about 10° off forward. Thus the Compton electrons are concentrated near the forward direction.

The velocity v of a 1-MeV electron is such that

$$\frac{v}{c} \equiv \beta \approx 0.94 \quad , \qquad (12)$$

and the mass is about 3 times its rest mass. In the geomagnetic field appropriate for the central U.S.,

$$B_0 = 0.56 \text{ Gauss} \quad , \qquad (13)$$

the gyro (or Larmor) radius of such an electron is

$$L = 85 \text{ m} \quad . \qquad (14)$$

The mean stopping range of the 1-MeV electron, due to collisions with electrons in air atoms, at 30 km altitude, is

$$R_m = 170 \text{ m} \quad . \qquad (15)$$

This range has been reduced from the extreme range (R_e = 245 m) to account for the multiple scattering of the electron by the nuclei of air atoms. The energy lost by the 1-MeV electron in stopping produces about 30,000 secondary electrons distributed along its path. The secondary electrons have kinetic energies of the order 10 eV at birth. Because their velocities are randomly distributed, they make no significant current of their own accord. However, if an electric field is present, they drift in the direction of the electric force on them, forming an electric conductivity in the air. The conductivity is discussed further in Section 2.4.

Because the mean range of the Compton electron is only about twice its gyro radius, it turns through only about 2 radians or 115° in the geomagnetic field before stopping. Because its mass and gyro radius decrease as it loses energy, it actually turns through a larger angle than this. The average trajectory of the Compton electrons at 30 km altitude will look like that sketched in figure 1.

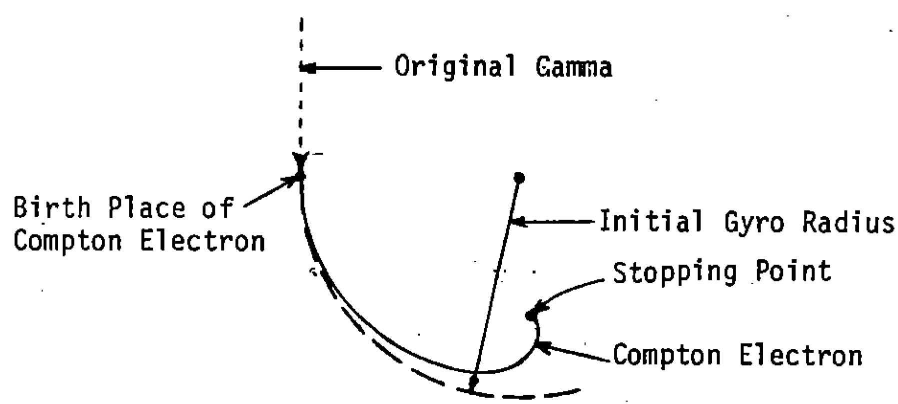

Figure 1. Sketch of average trajectory of Compton electrons from 2-MeV gammas at 30 km altitude.

We can now understand the difference between the density of birth places of Compton electrons and the density of moving Compton electrons. Let us have an impulse function of gammas moving through air. The trajectory of this gamma pulse, i.e., its position r as a function of

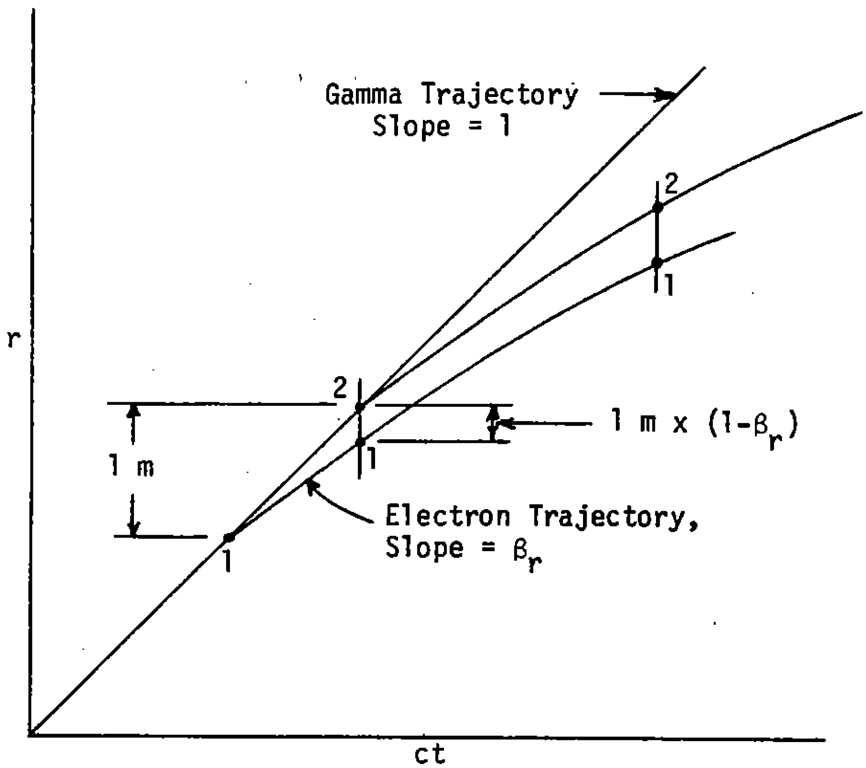

Figure 2. Explanation of motion compression of Compton electrons.

ct (t is time) is indicated in figure 2. Consider also two Compton electrons with birthplaces 1 m apart in r. The trajectories of these electrons, which start with velocity $v_r/c \equiv \beta_r$ slightly less than unity, are identical except that trajectory 2 is translated parallel to the gamma trajectory with respect to that of trajectory 1. The curvature of the electron trajectories is exaggerated in the figure. As indicated in the figure, at the time when the second electron is born, the first electron is a distance 1 m × (1-β_r) away from the second. Since all electrons born between 1 and 2 will be in this reduced interval, the relation between the density N_{ce} of Compton electrons and the density N_{bp} of birthplaces is

$$N_{ce} = \bar{N}_{bp}/(1-\beta_r) \quad . \tag{16}$$

It is not difficult to show that this same relation holds as β_r decreases due to geomagnetic deflection, energy loss and scattering. The factor $1/(1-\beta_r)$ must, of course, be averaged over the angular distribution of the Compton electrons. For 2-MeV gammas, this average is

$$\text{av}\left(\frac{1}{1-\beta_r}\right) \approx 11.2 \tag{17}$$

just after birth of the Compton electrons (Ref. 5). This justifies the factor of 10 used in Section 2.1. In calculations of the HEMP using the continuum Compton current, the expression for this current density is

$$\vec{J}_c = - \bar{N}_{bp} ec \ \text{av}\left(\frac{\vec{\beta}}{1-\beta_r}\right) \quad , \tag{18}$$

where -e is the electron charge. (It is actually the average of $\beta_r/(1-\beta_r)$ that has the value given in equation (17), but β_r is only slightly less than unity.)

2.3 FIELDS RADIATED BY MOVING ELECTRONS

We have seen that the Compton current density averaged over volumes of the order of 1 m^3 has only very small fluctuations in our nominal HEMP problem. This averaged current density is commonly called the macroscopic current density, and is the quantity used in almost all of the applications of electrical engineering, e.g., analysis of radio broadcasting and receiving, electric power generation and transmission, etc. Engineers (and physicists) never try to analyze these problems in terms of the fields radiated by individual electrons, but have nevertheless been getting correct answers for over 100 years (in fact, since before electrons were discovered). We therefore regard it as curious that quite a few people have insisted that the HEMP must be derived from the radiation of individual electrons in order to obtain reliable results. This seems all the more curious when one recalls that the fields of moving and accelerating electrons are derived in the textbooks by solving Maxwell's (continuum) equations for point charges and currents. Because those equations are linear, it must be true that the average or macroscopic fields are obtained by solving Maxwell's equations with the macroscopic currents.

The reason for the common use of the macroscopic current is that it is generally much easier to add up the currents of a large number of electrons than to add up their fields. However, with sufficient care the latter can be done correctly, as the following simple but relevant examples will show. In this section the essential physics of the examples is described, while the supporting mathematics is worked out in Appendix A.

Example 1. Impulse of Gammas on Thin Sheet

Let us have a planar, impulse function of gammas propagating through vacuum and arriving broadside on a thin insulating sheet, as indicated in figure 3. We imagine that electrons are knocked out of the sheet in the forward direction of the gammas, and that there is a magnetic field B_0 parallel to the sheet, in the y-direction. This field deflects the electrons in semicircles, until they restrike the sheet and are stopped. An observer is located at a distance z_0 below the sheet, which is very large compared with the gyro-radius of the electrons. (This is for ease of calculation.)

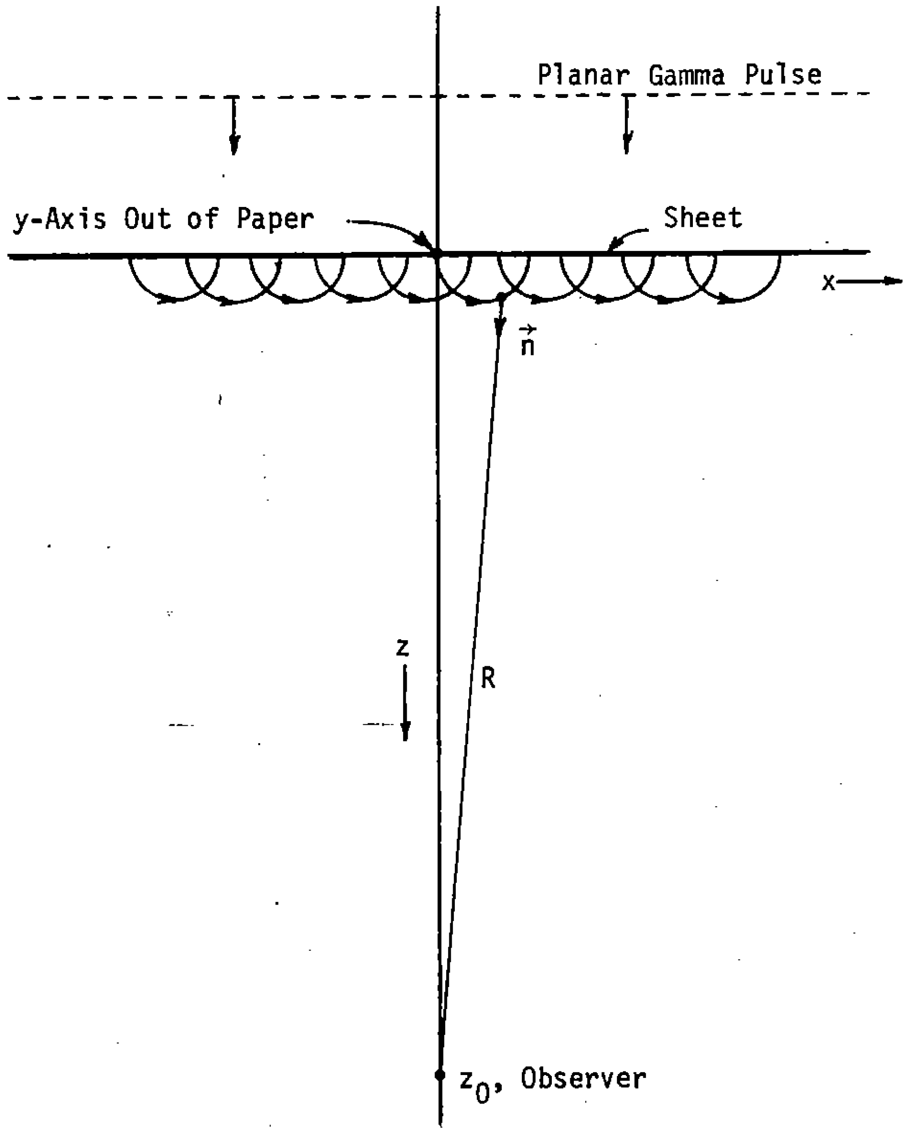

Figure 3. Configuration for Example 1.

The power flow density in the wave is

$$P(T,z) = E_x^2(T,z)/Z_0 \text{ watts/m}^2 \quad , \tag{60}$$

where

$$Z_0 \approx 120\pi \approx 377 \text{ ohms} \quad . \tag{61}$$

With sources and absorbers present, observers at different z may see different power densities at the same T. Energy is put into or taken from electromagnetic fields by currents flowing against or in the direction of the electric field. If $J_x(T,z)$ is the current density in Amps/m^2, then

$$J_x E_x \text{ watts/m}^3 \tag{62}$$

is the power/m^3 being taken out of the fields and put into the current (electrons). This energy transfer to the electrons must decrease the power flow density in the wave. Considering two z's a differential dz apart, we can write the law of energy conservation as

$$\frac{\partial}{\partial z} \frac{E_x^2}{Z_0} \left(\frac{\text{watts}}{m^3}\right) = - J_x E_x \left(\frac{\text{watts}}{m^3}\right) \quad . \tag{63}$$

The partial derivative here means that T is held constant; energy taken out of the wave near time T reduces the power flow density near that time, and the notch in the wave so produced propagates along with the wave. Now carrying out the differentiation of E_x^2 gives

$$\frac{\partial}{\partial z} E_x = - \frac{Z_0}{2} J_x \quad . \tag{64}$$

Here J_x is the total current density. If we have both Compton current J_{xc} and conduction current σE_x, with σ the conductivity, then

$$\frac{\partial}{\partial z} E_x = -\frac{Z_0}{2}(J_{xc}+\sigma E_x) \quad . \tag{65}$$

This is the outgoing wave equation in planar geometry, first derived in Reference 1. It is called outgoing because the wave propagates in the same direction as the gammas. The Compton current, driven by the gammas, appears therefore to move with the gammas, so that the variables T,z are also convenient for expressing $J_{xc}(T,z)$.

Application to Example 3

In Example 3 we neglected any conductivity. Putting $\sigma = 0$ in equation (65) and integrating z through the source region gives

$$E_x(T) = -\frac{Z_0}{2}\int J_{xc}(T,z)dz \quad . \tag{66}$$

The macroscopic Compton current density is

$$J_{xc} = -ev_x N_{ce} \quad , \tag{67}$$

where $v_x = \beta_x c$ is the electron velocity component and N_{ce} is the density of Compton electrons, which is related to the density N_{bp} of Compton electron birthplaces by equation (16). Thus

$$E_x(T) = \frac{Z_0}{2} ec \int \frac{N_{bp}(z)\ \beta_x(T)}{1-\beta_z(T)} dz \quad . \tag{68}$$

For Example 3, β_x and β_z are given by equation (22), i.e.,

$$\beta_z = \beta\cos\omega_e t \quad , \quad \beta_x = \beta\sin\omega_e t \quad , \tag{69}$$

where t is related to T by (equation (A-7))

$$\omega_e T = \omega_e t - \beta\sin\omega_e t \quad . \tag{70}$$

Therefore the final expression for E_x is

$$E_x(T) = \frac{Z_0}{2} e\beta c \left(\int n_{bp} dz\right) \frac{\sin\omega_e t}{1-\beta\cos\omega_e t} \quad . \tag{71}$$

This is identical with equation (50), with equation (53) for N_A. Recall that N_A is also equal to the number of gammas per m^2 incident on the atmosphere.

Probably most readers will agree that the calculation of E_x here is simpler and more reliable than that in Section 2.3 based on adding the fields radiated by individual electrons.

The peak value of E_x occurs at $\cos\omega_e t = \beta$ and is

$$E_{xp} = \frac{Z_0}{2} \frac{e\beta c}{\sqrt{1-\beta^2}} N_A \quad . \tag{72}$$

With $\beta = 0.94$ (1-MeV electrons) and $N_A = 6.5\times10^{-14}/m^3$ (equation (9)), this becomes

$$E_{xp} = 1.6\times10^7 \text{ V/m} \quad . \tag{73}$$

The HEMP is actually never this large, the chief reason being the severe attenuation due to air conductivity. While the effect of air conductivity is easily calculated from the outgoing wave equation, it is very difficult to include accurately in the individual electron calculation.

Spherical Geometry

The derivation of the outgoing wave equation given above for planar geometry is easily modified for spherical geometry. In this case the gammas come from a source with dimensions of the order of a meter in a time span of the order of 10^{-8} second. Thus at later times the gammas occupy a spherical shell, centered about the burst point, and with thickness of a few meters. The Compton current is produced in the part of this shell that intersects the atmosphere at altitudes below 50 km or so. It is clear that outgoing spherical waves will be generated if the Compton current has components transverse to the radius vector from the burst point, and the geomagnetic field ensures that there will be such currents. The wavelengths in these waves will be of the order a few meters to tens of meters (to fit into the shell). Because of the large radius of the shell, from burst point to source region, variation of the wave with transverse distances will be very slow compared with variation with radial distance. This means that the waves propagate approximately in the radial direction, since they propagate in the direction of their gradients.

If observers trigger their scopes on arrival of the gamma pulse, then the delayed time T for them in the spherical case is related to standard time t by

$$T = t - \frac{r}{c} \quad , \tag{74}$$

where r is the distance from the burst point. The transverse wave electric field E_t will now be a function of T and r. In a fixed element $\delta\Omega$ of solid angle, the power flow is

$$P(T,r) = \frac{E_t^2(T,r)}{Z_0} r^2 \delta\Omega \text{ watts} \quad . \tag{75}$$

If there is a current J_t in the direction of E_t, then the power flow will vary with r according to

$$\frac{\partial}{\partial r}\left(\frac{E_t^2(T,r)}{Z_0} r^2 \delta\Omega\right) = - J_t(T,r) E_t r^2 \delta\Omega \frac{\text{watts}}{\text{m}} \quad ,$$

or

$$\frac{1}{Z_0}\frac{\partial}{\partial r}(rE_t)^2 = r^2 J_t E_t \quad . \tag{76}$$

Carrying out the derivative leads to

$$\frac{\partial}{\partial r}(rE_t) = -\frac{Z_0}{2} rJ_t$$

$$= -\frac{Z_0}{2} r(J_{tc}+\sigma E_t) \quad . \tag{77}$$

Here we have expressed the total transverse current as the sum of Compton and conduction parts. This is the outgoing wave equation in spherical geometry, and was first derived by Karzas and Latter in Reference 2.

Comparison of equations (65) and (77) shows that the only difference is the replacement of E_t by rE_t in the spherical case and J_t by rJ_t. This replacement coincides with the fact that, in a case without Compton current and conductivity, E_t is independent of z in the planar case while rE_t is independent of r in the spherical case.

If the burst point is very far away from the source region, then the factor r varies little in the source region. In this case, the planar equation (65) can be used in the source region, but we should remember that the amplititude of the HEMP will fall, after it is produced, inversely proportional to the distance from the burst point.

Equations (65) and (77) replace the full set of 6 Maxwell equations with partial derivatives in 3 space coordinates and time. They make it possible to understand the HEMP in quite simple terms, and to find approximate analytical solutions. They also make it possible to find numerical solutions, that are accurate to a few percent, with modest rather than prohibitive amounts of computer time.

Effect of Conductivity; Saturation

It is easy to understand the effect of conductivity from either of equations (65) or (77). In the planar case, if there is no Compton current, $J_{xc} = 0$, then E_x attenuates with distance according to

$$E_x(T,z) = E_{xo}(T) \exp \left[- \frac{Z_0}{2} \sigma z\right] \quad . \qquad (78)$$

The attenuation length is $2/Z_0\sigma$, which becomes shorter for larger σ. Thus two effects oppose each other in equation (65). As the wave moves along in z, its amplitude is increased by the Compton current in the direction of $-J_{xc}$. The minus sign here is an example of Lenz's law. At the same time, the wave is attenuated by the conductivity. These two opposing effects balance if the right-hand side of equation (65) vanishes, i.e., $\partial E_x/\partial z = 0$ if

$$E_x = - J_{xc}/\sigma \quad . \tag{79}$$

This field, for which the Compton current is cancelled by the conduction current, is called the saturated field. The HEMP field approaches this value in the source region when the attenuation length becomes small compared with the thickness of the source region. Note that the saturated field is the same in planar and spherical geometry.

Below the source region, where the gammas have been mostly scattered, both J_{xc} and σ become small, but the saturated field changes only a little. When J_{xc} and σ become sufficiently small, the HEMP propagates as a free wave without further buildup or attenuation. The 1/r dependence is still present, of course.

The solution of the outgoing wave equation is discussed in detail in References 5 and 6.

Diffraction

In our derivation of the outgoing wave equations, we assumed that energy flow was strictly in the z- or r- directions in the planar or spherical cases, respectively. In the real, spherical case, this was justified by the argument that the transverse variations of the wave are small compared with the radial variations. The wave must have transverse variations since there are no spherical waves that are independent of the angles of spherical coordinates. The wave must be composed of spherical harmonics Y_ℓ^m (in the standard notation) with $\ell \geq 1$. Generally, however, the minimum length scale of the transverse variations is set by the atmospheric scale height, equation (8), since the gamma flux and Compton current will vary significantly in this distance. For observers that are near the horizon from a given burst point, as sketched in figure 5, the vertical direction is not far from the θ-direction of spherical coordinates. In this case, the scale length of wave variations in θ is indeed approximately the scale height h. The scale length for variations in the azimuthal direction ϕ is comparable with the radius of the source region from the burst point, which is typically much greater than h; this variation comes from the changing angle between the newborn Compton electrons and the geomagnetic field. Thus the principal transverse variations are those associated with the scale height.

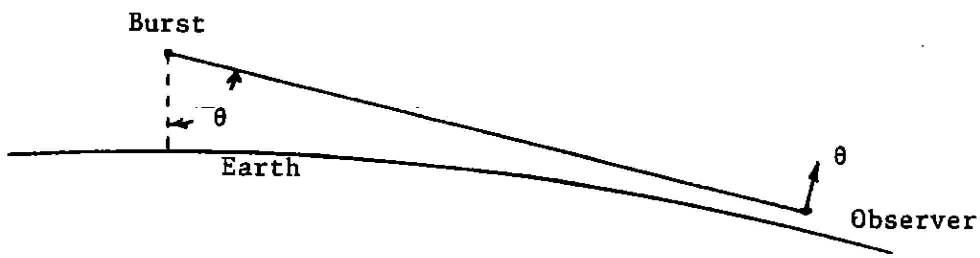

Figure 5. For an observer near the horizon, the θ-direction is approximately vertical.

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